Production of large-area plasmas by electron beams*

R. F. Fernsler,[†,a)] W. M. Manheimer, R. A. Meger, J. Mathew, D. P. Murphy, R. E. Pechacek,[b)] and J. A. Gregor[c)]
Naval Research Laboratory, Beam Physics Branch, Code 6750, 4555 Overlook Ave., S.W., Washington DC 20375-5346

(Received 19 November 1997; accepted 10 February 1998)

An analysis is presented for the production of weakly ionized plasmas by electron beams, with an emphasis on the production of broad, planar plasmas capable of reflecting X-band microwaves. Considered first in the analysis is the ability of weakly ionized plasmas to absorb, emit and reflect electromagnetic radiation. Following that is a determination of the electron beam parameters needed to produce plasmas, based on considerations of beam ionization, range, and stability. The results of the analysis are then compared with a series of experiments performed using a sheet electron beam to produce plasmas up to 0.6 m square by 2 cm thick. The electron beam in the experiments was generated using a long hollow-cathode discharge operating in an enhanced-glow mode. That mode has only recently been recognized, and a brief analysis of it is given for completeness. The conclusion of the study is that electron beams can produce large-area, planar plasmas with high efficiency, minimal gas heating, low electron temperature, high uniformity, and high microwave reflectivity, as compared with plasmas produced by other sources. © *1998 American Institute of Physics.* [S1070-664X(98)95505-7]

I. INTRODUCTION

Nonequilibrium, weakly ionized plasmas have a number of unique properties that can be exploited for a wide variety of industrial and scientific applications. One such property is the ability to decay rapidly once the ionization source is removed, and a second property is the ability to reflect electromagnetic waves below the plasma frequency. As suggested by Manheimer,[1] a thin planar plasma could thus serve as a lightweight agile mirror used, for example, to direct a radar beam to and from a fixed receiver/source. Unlike metallic reflectors, weakly ionized plasmas can be created and destroyed quickly to provide rapid reorientation of the radar. A large, planar plasma could thus provide the directional capability of a phased-array antenna, but with a broad instantaneous bandwidth and potentially low cost.

A third useful property of nonequilibrium plasmas is the ability to produce large numbers of chemically reactive neutrals from a cold background gas. The reactive neutrals are created as free electrons excite and dissociate the gas molecules. Because the electrons are hot and energetic, the production rates are large even though the gas is cold. High degrees of gas excitation, dissociation, and ionization can thus be produced and maintained in gases at near room temperature. As a result, materials adjoining the plasmas can be processed quickly and with little thermal stress, and nonequilibrium plasmas are now widely utilized in thin-film processing of integrated circuits and other surface treatments.[2]

For the past several years the Naval Research Laboratory (NRL) has been producing and studying[3–5] planar plasmas up to 0.6 m square by 2 cm thick, as a first step toward developing an agile mirror for radar. Plasmas with densities above 10^{12} cm^{-3} are routinely produced and have been shown reflect 10 GHz microwaves as well as does a metallic plate.[3] Similar or lower densities are needed for plasma processing.[2] However, in contrast to the sources used in plasma processing, the present source produces a plasma that is thin yet uniform over a large area. This is advantageous for applications requiring large surfaces and/or high throughput. Moreover, the efficiency of plasma production is potentially higher than that of other sources, in terms of both energy and geometry, while the electron temperature is lower and independently controllable. Broad, dense, planar plasmas can thus be produced efficiently and with high control.

The source of ionization in the present experiments is an electron beam produced by a long hollow-cathode discharge. The discharge can operate in one of several modes, but the one that produces the highest overall plasma density operates in a mode dubbed[3] an ''enhanced glow.'' This mode occurs only for certain cathode configurations, and has been produced using cathodes ranging from 15 to 60 cm long. The length is limited only by the size of the vacuum chamber, and curved cathodes could be used if desired. The cathodes that produce an enhanced glow discharge are hollow with an internal cross section of $1-2$ cm^2 and an opening ~ 1 cm wide on one side; see Fig. 1. The open side faces a flat circular anode situated up to 60 cm away. At gas pressures between 20 and 300 mTorr, a voltage of 2–5 kV is impressed across the electrodes. Depending on the discharge parameters, an enhanced glow is produced, and an electron beam emerges with nearly the full discharge voltage and a current of $10-50$ mA/cm^2. An axial magnetic field ~ 200 G is used to confine the beam. At higher pressures or lower

*Paper oThaI2-5, Bull. Am. Phys. Soc. **42**, 2006 (1997).
[†]Invited speaker.
[a)]Electronic mail: fern@ppd.nrl.navy.mil
[b)]SFA, Inc., Landover, Maryland 20785.
[c)]University of Maryland, College Park, Maryland 20742

 © 1998 American Institute of Physics

Downloaded 18 Dec 2006 to 132.250.158.37. Redistribution subject to AIP license or copyright, see http://pop.aip.org/pop/copyright.jsp

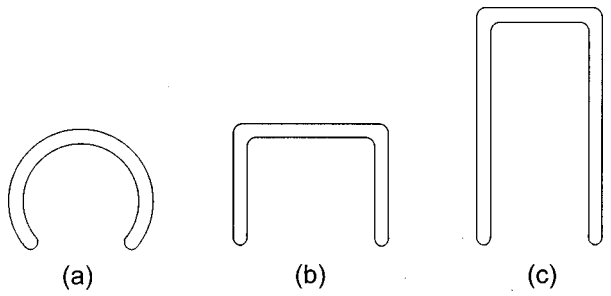

FIG. 1. Cross-sectional view of cathodes used in the NRL experiments: (a) a pipe, with inner diameter 1.3 cm and a slot 1 cm wide; (b) a shallow rectangle, 1.2 cm deep and 1.6 cm wide; and (c) a deep rectangle, 2.4 cm deep and 1.3 cm wide. Only the first two configurations are capable of producing an enhanced-glow discharge.

voltages, the beam is unable to propagate all the way to the anode, and a positive column forms from the end of the beam range to the anode; see Fig. 2. The plasma density in the positive column is much lower, however, and thus this mode is of little interest here. Other discharge modes exist as well.

To assess the plasmas being produced, we present a three-part analysis. The first part considers the plasma properties needed to reflect 10 GHz microwaves. The second part determines the electron beam parameters required to produce such a plasma. And the third part models the beam source presently used, namely an enhanced glow discharge. A more detailed discussion of the experiments themselves may be found in Refs. 3–5.

II. PLASMA REQUIREMENTS

The response of a weakly ionized, magnetized plasma to an incident electromagnetic wave depends on the angular frequency ω of the wave, the plasma frequency $\omega_{pe}=(4\pi e^2 n_{pe}/m)^{1/2}$, the electron cyclotron frequency $\omega_c=eB/mc$, and the momentum-transfer collision frequency ν_{pe} of the plasma electrons. Here n_{pe} is the plasma electron density, e is the electron charge, m is the electron mass, c is the speed of light, and B is the magnetic field strength. As shown in Appendix A, the wave is entirely reflected or absorbed if the plasma is overdense ($\omega_{pe}>\omega,\omega_c$) for several electromagnetic skin depths ($x>3c/\omega_{pe}$). Furthermore, if $\omega_{pe}>1.1\omega\gg\nu_{pe},\omega_c$, the fraction of wave energy absorbed is limited to

$$f_p<\frac{3}{2}\frac{\nu_{pe}}{\omega_{pe}}+\frac{8}{3}\frac{\nu_{pe}\Delta x}{c}. \tag{1}$$

The first term on the right represents the energy absorbed in the overdense plasma middle where $\omega_{pe}(x)>\omega$, while the second term represents the energy absorbed in the underdense wings where $\omega_{pe}(x)\leqslant\omega$. The width of the wings is Δx.

High reflectivity requires low absorption, $f_p\ll 1$. According to Eq. (1), low absorption requires a large peak plasma frequency ω_{pe}, a small collision frequency ν_{pe}, and narrow wings Δx. The plasma frequency is determined by the electron density n_{pe}, while the collision frequency is determined by the collision cross section, the gas density N, and the electron temperature T_e. In gases like oxygen,[6] $\nu_{pe}/N\cong 5\times10^{-8}$ cm^3/s for $T_e\sim 1$ eV. To achieve a reflectivity of 95% or better ($f_p\leqslant 0.05$) at 10 GHz thus requires (i) a dense plasma, $n_{pe}>1.2\times10^{12}$ cm^{-3}, for a cm or more, to keep $\omega_{pe}>\omega\cong 6\times10^{10}$ rad/s; (ii) a low gas pressure, typically below 300 mTorr, to keep $\nu_{pe}<5\times10^8$ s^{-1}; and (iii) narrow plasma wings, $\Delta x<1$ cm. In addition, the plasma length and width should exceed ten wavelengths, $\lambda=2\pi c/\omega=3$ cm, to minimize diffraction, and the plasma should be smooth and stationary to within 0.1λ to maintain phase coherence. Since the plasma screens out spatial fluctuations to a few Debye lengths, $\lambda_D=(T_e/4\pi e^2 n_{pe})^{1/2}\lesssim 10^{-3}$ cm, phase coherence is determined mainly by the smoothness of the plasma source.

High reflectivity has another benefit, namely low microwave emission from the plasma. Although unimportant for most applications, low emission is critical when detecting weak radar return signals. In the frequency range of interest, the effective plasma radiation temperature is related to the electron temperature through Kirchoff's law,

$$T_r\approx f_p T_e. \tag{2}$$

The radiation temperature can therefore be low, $T_r\lesssim 600$ K for $f_p\lesssim 0.05$, even though the plasma electrons are hot, T_e

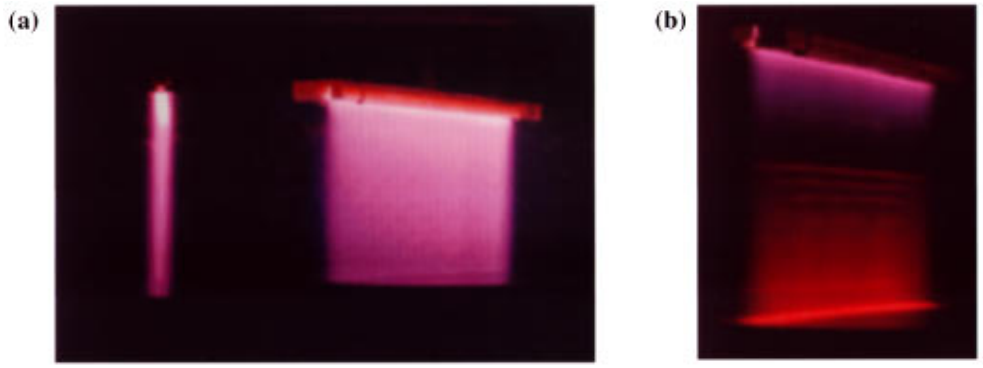

FIG. 2. Glow discharges produced using the rectangular cathode shown in Fig. 1(b). The cathode is 60 cm long and lies 60 cm away from a flat anode. An enhanced-glow discharge (a) is shown end-on and at an angle, in 120 mTorr of air at 4 kV and 250 G. The formation of a positive column (the orange–red region) is seen (b) when the pressure is increased to 200 mTorr and the voltage is decreased to 3 kV. The latter case exhibits a Faraday dark space and striations.

Downloaded 18 Dec 2006 to 132.250.158.37. Redistribution subject to AIP license or copyright, see http://pop.aip.org/pop/copyright.jsp

$\cong 1$ eV. Experimentally, the emission is small above the peak plasma frequency, but not below it. The source of this additional emission is discussed later.

III. BEAM REQUIREMENTS

An electron beam must meet four general criteria to generate the plasma described above. First the beam must have a range and width more than $10\lambda \geqslant 30$ cm. Second, the beam should remain confined to a thin sheet. Third, the beam must have enough current to produce the ionization needed. And fourth, the beam should be stable.

The beam range is limited both by elastic and inelastic scattering off neutrals. Inelastic scattering limits the range to

$$R \leqslant R_i \approx \frac{\varepsilon_b}{2(d\varepsilon_b/dz)}, \tag{3}$$

where the energy loss function $d\varepsilon_b/dz$ varies[7] as the gas density N times $\ln(\varepsilon_b/\varepsilon_b)$ for electron energies $\varepsilon_b \geqslant 500$ eV. The range R_i thus varies as $\varepsilon_b^2/[N \ln(\varepsilon_b)]$. As an example, consider 2 keV electrons propagating in 30 mTorr ($N \cong 10^{15}$ cm^{-3}) of oxygen at room temperature. Then[8] $d\varepsilon_b/dz \cong 3$ eV/cm, so that $R_i \cong 3$ m. However, if the electrons acquire appreciable transverse energy, the axial energy loss function $d\varepsilon_b/dz$ increases while the axial range R_i decreases, due to the longer pathlength. Nevertheless, axial ranges $R_i \sim 1$ m are readily obtainable for energies of several keV and gas pressures below 300 mTorr. Increasing the range to beyond the device length improves plasma uniformity but wastes energy.

Elastic gas scattering further limits the beam range, by converting forward energy into transverse energy:

$$\begin{aligned}\frac{d}{dz}\left(\frac{mv_\perp^2}{2}\right) &= -\frac{d}{dz}\left(\frac{mv_b^2}{2}\right) \\ &= -v_b\frac{d}{dz}(mv_b) = -\frac{d}{dt}(mv_b) \\ &= \nu_b m v_b = N\sigma_m m v_b^2.\end{aligned} \tag{4}$$

Here $v_\perp$ is the transverse beam velocity, v_b is the forward velocity, $\nu_b = v_b N\sigma_m$ is the collision rate, and σ_m is the momentum-transfer cross section. The reduction in v_b limits the axial propagation distance to

$$R \leqslant R_e \approx (N\sigma_m)^{-1}. \tag{5}$$

For gases like oxygen, the elastic range R_e is comparable to the inelastic range R_i, and it has a similar dependence[7] on beam energy ε_b and gas density N.

The increase in $v_\perp$ causes the beam to spread, but the spreading can be reduced by applying a strong magnetic field B. If $\omega_c = eB/mc \gg \nu_b$, the increase in $v_\perp$ converts immediately to an increase in the electron gyroradius, $r_c = v_\perp/\omega_c$, so that the mean-squared beam radius grows as

$$\frac{d}{dz}r_b^2 = \frac{d}{dz}\langle r_c^2\rangle = \omega_c^{-2}\left\langle\frac{dv_\perp^2}{dz}\right\rangle = 2N\sigma_m\left(\frac{v_b}{\omega_c}\right)^2 \approx \frac{2}{R_e}r_{\rm cm}^2, \tag{6}$$

where $r_{\rm cm} = v_b/\omega_c$ is the maximum gyroradius. A strong field thus keeps the expansion under $\sqrt{2}r_{\rm cm}$ for ranges $R < R_e$. To keep the plasma wings narrow, $r_{\rm cm}$ should be a cm or less. This requires $B \geqslant 200$ G for 2 keV electrons. The electron cyclotron frequency, $\omega_c \cong 3.5\times 10^9$ s^{-1}, is then small compared with ω_{pe} but large compared with ν_b.

A third requirement is to produce enough ionization to offset losses to recombination, attachment, and diffusion. For the agile-mirror application, the loss processes must, in addition, destroy the plasma rapidly once the ionizing beam turns off. Attachment is generally too slow at gas pressures of interest, while diffusion is destructive only at the walls. Dissociative recombination is both fast and destructive, $e + AB^+ \rightarrow A + B$, and it is generally the dominant loss mechanism in weakly ionized molecular gases. Molecular gases are commonly used in plasma processing applications as well,[2] because of the need for free radicals like Cl and F.

To sustain a plasma of density n_{pe} in a molecular gas, the beam ionization rate must satisfy

$$S_b = \frac{j_b}{e}\frac{1}{w_i}\frac{d\varepsilon_b}{dz} = \beta n_{pe}^2, \tag{7}$$

where j_b is the axial beam current density, w_i is the mean energy deposited by the beam per plasma electron created,[9] and β is the dissociative recombination rate.[10] For energetic beams, w_i is constant while $d\varepsilon_b/dz \cong \varepsilon_b/2R$, according to Eq. (3). Equation (7) then indicates that j_b varies as $n_{pe}^2 R/\varepsilon_b$ and that the beam power, $j_b\varepsilon_b$, varies as $n_{pe}^2 R$; these relationships remain valid even when the electrons gyrate under the influence of B, although the coefficients of proportionality change. For a given range R, the beam current and power thus increase as $n_{pe}^2 \propto \omega_{pe}^4$, and therefore a *16-fold* increase in beam current and power is needed to double the plasma frequency. To illustrate, consider again 2 keV electrons propagating in 30 mTorr of oxygen. Then $d\varepsilon_b/dz \cong 3$ eV/cm, while $w_i \cong 31$ eV and $\beta \cong 2\times 10^{-8}$ cm^3/s for $T_e \sim 1$ eV. Inserting these values into Eq. (7), we find that $j_b \cong 50$ mA/cm^2 is needed to produce $n_{pe} = 1.2\times 10^{12}$ cm^{-3} (10 GHz reflection), but that $j_b \cong 800$ mA/cm^2 is needed to produce $n_{pe} = 5\times 10^{12}$ cm^{-3} (20 GHz reflection). The maximum plasma frequency achievable is thus severely limited by the available beam flux. Once the beam turns off, the plasma decays in a time $\tau_d = (\beta n_{pe})^{-1}$, which ranges from 10 to 40 μs in the examples given. Note that for applications requiring very dense plasmas but *not* rapid decay, an atomic gas like argon should be considered instead, since β is much smaller for atomic ions.

The fourth and final beam requirement is stability. At the current densities indicated, beam self-forces are negligible, and the beam electrons respond mainly to the applied magnetic field, collisions with gas molecules, and electromagnetic interactions with the plasma. The latter can give rise to instability, but only if the instability grows faster than the collisional damping rate ν_{pe}. One of the few instabilities capable of growing faster than ν_{pe} is the two-stream instability. This instability is detuned, however, by variations in the axial beam velocity v_b, and it is fully suppressed if the variations are large enough.[11] As shown in Appendix B, if

Downloaded 18 Dec 2006 to 132.250.158.37. Redistribution subject to AIP license or copyright, see http://pop.aip.org/pop/copyright.jsp

the velocity distribution falls linearly to zero over a range Δv from a maximum velocity v_0, the beam is stable provided the plasma electron collision frequency satisfies

$$\nu_{pe} \geqslant 2\pi\omega_{pe}\left(\frac{n_b}{n_{pe}}\right)\left(\frac{v_0}{\Delta v}\right)^2. \tag{8}$$

Here $n_b = j_b/ev_m \ll n_{pe}$ is the beam density, and $v_m = (v_0 - \Delta v/2)$ is the mean beam velocity. Since $\Delta v \leqslant v_0$, the minimum stability criterion for a linear velocity spread is

$$\nu_{pe} > 2\pi\omega_{pe}\left(\frac{n_b}{n_{pe}}\right). \tag{9}$$

To understand the significance of condition (9), recall that j_b, and therefore n_b, is proportional to n_{pe}^2. As a consequence, the right-hand side of Eq. (9) increases as $(n_{pe})^{3/2} \propto \omega_{pe}^3$. Since the collision frequency ν_{pe} is determined mainly by the gas pressure, the stability criterion cannot be satisfied above some minimum ω_{pe} determined by the pressure. In the two examples discussed earlier, the first case ($n_{pe} = 1.2 \times 10^{12}$ cm^{-3}) is stable, provided $\Delta v \gtrsim v_0/2$, while the second case ($n_{pe} = 5 \times 10^{12}$ cm^{-3}) is unstable for all Δv. The two-stream instability therefore requires not only a large velocity spread, but it limits the maximum plasma frequency achievable. This limit can be raised by increasing the gas pressure to increase ν_{pe}, but at the cost of higher bremsstrahlung emission. Of course, the two-stream instability itself produces radiation, mainly at the local plasma frequency. Such emission is consistent with the noise spectrum observed, which indicates that the discharges described in Refs. 3–5 are only marginally stable.

IV. ADVANTAGES OF BEAM-PRODUCED PLASMAS

In the previous section we showed that electron beams can create dense plasmas provided the beam is stable, magnetically confined, and has sufficient current and energy. The plasma created by the beam should then be smooth axially, and smooth transversely if the beam is smooth. Moreover, if recombination is the dominant loss mechanism, the plasma produced will extend only slightly outside the beam. A sheet beam can thus produce a sharp-edged, planar plasma capable of reflecting microwaves with negligible loss.

Another advantage of electron beams is that they require less power than other plasma sources. To show this, consider the volumetric power, $w_i n_{pe}/\tau_d$, needed to sustain a plasma of density n_{pe}. Here w_i is the mean energy deposited in the gas per electron–ion pair created, and τ_d is the average lifetime of a plasma electron. Since n_{pe} and τ_d are essentially fixed for a given plasma, the power needed is determined by w_i alone. For sources that produce ionization by heating the plasma electrons, w_i is many times the ionization energy ε_i of the gas molecules.[2] This is because only a small fraction of the plasma electrons have enough energy to ionize, so that most of the plasma energy is absorbed by low-lying excitations rather than ionization. By contrast, all the electrons in an energetic beam can ionize, and in that case[9] $w_i < 3\varepsilon_i$. Because w_i is smaller for electron beams, beams require less power to produce a given plasma. By reducing the power deposited, heat removal and gas flow requirements are alleviated as well.

Another property of beam-produced plasmas is that most of the energy deposited ultimately ends up in gas dissociation. This is because the cross sections for dissociation are large, and because dissociative recombination converts molecular ions into atoms. Beams are thus efficient at dissociating, as well as ionizing, molecular gases.

Finally, the electron temperature in beam-produced plasmas is intrinsically lower than that of other sources. Although beams create plasma electrons with energies of several eV or more, inelastic collisions with molecules cool these electrons in a time $\tau_c \gtrsim 10/\nu_{pe}$. Molecular gases, in particular, keep $T_e < 1$ eV, due to vibrational cooling.[6] Should high T_e be *required*, the plasma can be heated through rf heating or other means. In the enhanced-glow discharge discussed later, for example, a dc longitudinal electric field keeps $T_e \cong 1$ eV. The ability to control T_e, independent of the ionization rate, provides a degree of control generally unavailable with other sources.

V. BEAM PRODUCTION

To produce the plasmas desired, a sheet electron beam is required with a current density of several tens of mA/cm^2, an energy of several keV, and a large spread in axial velocity. The beam must, in addition, be injected into gas at 20–300 mTorr pressure. One means of generating an electron beam in gas is a glow discharge, and here we show that glow discharges can produce the requisite beam only when operated in a special mode termed an enhanced glow.[3]

A dc glow discharge consists of a cold cathode and anode in gas.[2] The disruption of electron flow at the cathode leads to the formation of an ion-rich sheath, and this sheath restricts the current much like a Child–Langmuir sheath. However, unlike a vacuum sheath, the current in a glow discharge transitions from primarily ion current at the cathode to electron current in the bulk plasma, as may be seen from the electron and ion continuity equations in steady state:

$$\nabla \cdot j_i = eS_i = -\nabla \cdot j_e. \tag{10}$$

Here j_i is the ion current density, j_e is the electron current density, and S_i is the net, volumetric production rate for electrons and ions. Equation (10) shows that the total current is conserved, $j_i + j_e = (1+\gamma) j_{ic}$, where j_{ic} and γj_{ic} are the ion and electron current densities, respectively, at the cathode, and $\gamma \lesssim 0.3$ is the number of electrons emitted from the cathode per incident ion.[12]

Outside the sheath, the ion and electron densities are nearly equal, $n_i \cong n_e$, while $S_i \cong 0$ because of the competition between ionization and loss processes like electron–ion recombination. In this region, j_i and j_e are essentially constant. However, the fluid velocity of the massive ions is small compared with that of the electrons, $v_i \ll v_e$, and therefore $j_i = en_i v_i \ll j_e = en_e v_e \cong (1+\gamma) j_{ic}$.

The ion current density in a glow discharge thus rises from a small fraction of $(1+\gamma) j_{ic}$ in the bulk plasma to j_{ic} at the cathode. According to Eq. (10), such a transition is

Downloaded 18 Dec 2006 to 132.250.158.37. Redistribution subject to AIP license or copyright, see http://pop.aip.org/pop/copyright.jsp

possible only if appreciable ionization occurs in the sheath. If beam-like electrons are to provide this ionization and pass N_r times through the sheath, the thickness of the sheath must satisfy

$$s \gtrsim \frac{\lambda_{ib}}{N_r}, \tag{11}$$

where λ_{ib} is the electron ionization mean-free path. Note that recombination is negligible in the sheath, because the electrons there are energetic[10] and few in number.

The electrostatic potential, ϕ, and field, $E=-\nabla\phi$, are governed by Poisson's equation,

$$\nabla^2\phi = 4\pi e(n_e - n_i) = -4\pi\left(\frac{j_e}{v_e}+\frac{j_i}{v_i}\right). \tag{12}$$

In the quasineutral bulk plasma $E\cong 0$, but in the non-neutral sheath E has been found[13,14] to rise linearly with distance x toward the cathode. As we now show, such behavior suggests that S_i is constant in the sheath (and zero in the bulk plasma).

The source S_i creates an ion flux, $dx_0\, S_i$, between x_0 and x_0+dx_0. At $x>x_0$, the ion density from this flux is given by $\delta n_i(x,x_0)=dx_0\, S_i/v_i(x,x_0)$, where $v_i(x,x_0)$ is the ion velocity. For collisionless ions created with zero energy at x_0, the velocity at x satisfies the energy constraint

$$\frac{Mv_i^2(x,x_0)}{2} = -e[\phi(x)-\phi(x_0)] = \frac{x^2-x_0^2}{s^2}\, e\phi_s, \tag{13}$$

where M is the ion mass. Here we have set $\phi(x)=-(x/s)^2\phi_s$, where ϕ_s is the voltage drop across the sheath. The total ion density at x equals the sum over all δn_i,

$$\begin{aligned} n_i(x) &= \int_0^x dx_0\, \frac{S_i}{v_i(x,x_0)} \\ &= \left(\frac{s}{x}\right)\int_0^x dx_0 S_i \sqrt{\frac{M}{2e\phi_s}}\left(1-\frac{x_0^2}{x^2}\right)^{-1/2} \\ &= \frac{\pi}{2}\sqrt{\frac{M}{2e\phi_s}}\, sS_i, \end{aligned} \tag{14}$$

assuming that S_i is constant and that the ion current entering the sheath is small, $j_i(0)\ll j_i(s)$. In that case $j_i(x)\cong exS_i$ and $j_i(s)\cong esS_i$, according to Eq. (10).

In the sheath $n_e\ll n_i$, because $j_e\lesssim j_i$ and $v_e\gg v_i$. Dropping n_e and setting $\nabla^2\phi=-2\phi_s/s^2$ in Eq. (12) and $sS_i=j_i(s)/e$ in Eq. (14), the ion current density at the cathode is given by

$$j_i(s)=j_{ic}\approx\sqrt{\frac{2e}{M}}\,\frac{\phi_s^{3/2}}{\pi^2 s^2}. \tag{15}$$

This solution yields values nearly three times larger than the vacuum Child–Langmuir solution.[15] Moreover, the solution is an upper bound because it neglects ion collisions which reduce v_i.

Depending on the cathode geometry, only a fraction k_e of the cathode-emitted electrons escape with full energy, $\varepsilon_b\cong e\phi_s$. The beam current density outside the cathode is thus $j_b=k_e\gamma j_{ic}$, with any excess current carried by the plasma. According to Eqs. (11) and (15), the beam current density is therefore restricted to

$$j_b \lesssim k_e\gamma N_r^2\sqrt{\frac{2e}{M}}\,\frac{\phi_s^{3/2}}{\pi^2\lambda_{ib}^2}. \tag{16}$$

In an abnormal glow discharge,[2] the cathode is effectively planar, and all of the cathode-emitted electrons escape after crossing the sheath once: $k_e=N_r=1$. In that case both the beam current density and the total current density, $j_t=(1+\gamma)j_{ic}$, are small. For example, for $\phi_s=2$ kV in 30 mTorr of oxygen, $j_b<j_t\lesssim 0.2$ mA/cm^2, assuming $\gamma\lesssim 0.3$ and[16] $\lambda_{ib}\sim 5$ cm. Cold, flat cathodes are thus generally incapable of producing large currents at modest voltages and pressures.

Much larger currents can be produced using hollow cathodes.[13] Electrons emitted from inner faces of the cathode then accelerate through the sheath toward opposing faces and reflect off. If the cathode hollow is narrow and deep, the electrons reflect many times and lose most of their energy before escaping: $N_r\gg 10$ and $k_e\to 0$. The total current density j_t is then large but j_b remains small. This is the usual operating mode of hollow-cathode discharges.[17]

However, if the cathode hollow is instead shallow with a large opening, electrons inside undergo only several reflections before escaping. An energetic beam of modest current then emerges from the hollow. This is the basis for the enhanced-glow cathode[3] depicted in Figs. 1(a) and 1(b). To keep the sheaths from overlapping, the dimensions of the hollow should satisfy $d>2s$. This constraint together with Eq. (15) limits the minimum current produced by an enhanced glow, and currents measured are consistent with that limit.[3,18]

To estimate the number of electron reflections, consider electrons emitted from an inner sidewall of a rectangular cathode like that shown in Fig. 1(b). Ignoring the magnetic field and assuming $s\ll d$, the electrons accelerate through the sheath and reflect back upon reaching the opposite side. At the same time, the electrons scatter off the gas molecules and thus slowly diffuse in the transverse direction. This diffusion is unaffected by specular reflection off the flat sheaths, and a spot of electrons therefore grows in radius as[19]

$$r(x)\approx(N\sigma_m x^3)^{1/2}, \tag{17}$$

where x is the total distance traveled and σ_m is the momentum-transfer cross section. Once $r\gtrsim d$, the electrons fill the hollow and begin to leave. This suggests that the electrons escape after crossing the two side sheaths a number of times, given by

$$N_r\approx 2x/d\approx 2(dN\sigma_m)^{-1/3}. \tag{18}$$

In the NRL experiments $d\cong 1.5$ cm, and therefore $N_r\sim 10$ for $\phi_s=2$ kV in 30 mTorr of oxygen, based[20] on $\sigma_m\cong 6\times 10^{-18}$ cm^2. Beam current densities as large as 20 mA/cm^2 are then possible from each sidewall. The total current density outside the hollow can therefore reach 40 mA/cm^2, or more, depending on the geometry of the cathode and the contribution from the back wall. Moreover, be-

Downloaded 18 Dec 2006 to 132.250.158.37. Redistribution subject to AIP license or copyright, see http://pop.aip.org/pop/copyright.jsp

cause electrons from the sidewalls leave with a large spread in longitudinal velocity, the two-stream instability is partially if not wholly suppressed. An enhanced-glow discharge is thus capable of producing moderately stable, high-current electron beams with energies of a few keV.

VI. SUMMARY

In this paper we have determined the properties needed for a sheet plasma to reflect microwaves efficiently up to 10 GHz, and then have determined the electron beam properties needed to produce such a plasma. The required plasma properties are (i) an electron density above 1.2×10^{12} cm^{-3} for a cm or more, and then falling to zero in a cm or less; (ii) a length and width more than 30 cm across; (iii) a surface that is stable and flat to 1 mm or less; and (iv) a gas pressure below 300 mTorr. We have shown, in addition, that such plasmas decay in 10–40 μs in molecular gases, due to recombination. The beam properties needed are (i) an energy of a few keV; (ii) a current density of several tens of mA/cm^2; (iii) a large spread in longitudinal velocity; (iv) magnetic confinement by fields $\sim$200 G; and (v) a gas pressure above 30 mTorr. Compared with other plasma sources, beam ionization was found to be more efficient, to minimize gas heating, and to produce cooler plasma electrons. And last, we have presented a qualitative model for the production of a sheet electron beam using an enhanced-glow discharge. Such discharges are predicted to occur when the cathode is hollow but shallow with a large opening on the side facing the anode. The analysis is in general agreement with recent experiments,[3–5] which have produced sheet plasmas up to 0.6 m square. The plasmas generated were stable for 1 ms or more, could reflect X-band microwaves as well as a metallic plate, and could be created and destroyed in times as short as 20 μs.

ACKNOWLEDGMENTS

The authors thank A. E. Robson for useful discussions and M. Myers for experimental support.

This work was supported by the Office of Naval Research.

APPENDIX A: ELECTROMAGNETIC ABSORPTION

Consider a uniform plasma subjected to a plane electromagnetic wave with an electric field $E(x,t)=E_y e^{i(kx-\omega t)}$. This field induces a plasma electron current density, given by

$$j_{pe}=\frac{e^2 n_{pe} E}{m(\nu_{pe}-i\omega)}, \tag{A1}$$

assuming the plasma electrons are cold and the dc magnetic field is weak. The field E is determined by Maxwell's equations,

$$\nabla x(\nabla xE)+\frac{1}{c^2}\frac{\partial^2 E}{\partial t^2}=-\frac{4\pi}{c}\frac{\partial j_{pe}}{\partial t}, \tag{A2}$$

where c is the speed of light. Combining these equations produces the dispersion relationship,

$$k^2c^2=\omega^2-\frac{\omega\omega_{pe}^2}{(\omega-i\nu_{pe})}, \tag{A3}$$

where $\omega_{pe}=(4\pi e^2 n_{pe}/m)^{1/2}$ is the plasma frequency. The neglect of the dc magnetic field is justified, provided the electron cyclotron frequency is small, $\omega_c=eB/mc\ll\omega$. This corresponds to magnetic fields $B<500$ G for frequencies of interest.

To analyze Eq. (A3), choose ω real and set $k=k_r+ik_i$, where k_r and k_i are the real and imaginary parts, respectively. Wave absorption is then small only if k_i is small, which requires

$$\nu_{pe}\ll\omega<\omega_{pe}. \tag{A4}$$

When $\nu_{pe}\ll\omega$, Eq. (A3) gives

$$k_i^2\approx\frac{\omega_{pe}^2-\omega^2-\nu_{pe}^2}{2c^2}+\sqrt{\left(\frac{\omega_{pe}^2-\omega^2-\nu_{pe}^2}{2c^2}\right)^2+\left(\frac{\nu_{pe}\omega_{pe}^2}{2\omega c^2}\right)^2}. \tag{A5}$$

The waves are evanescent if $\omega<\omega_{pe}$, and they effectively cease to propagate past $x=3c/\omega_{pe}$. There is thus little transmission through a plasma that is overdense for more than several electromagnetic skin depths. The incident wave is then reflected or absorbed, and the fraction of incident wave power, $cE^2/4\pi$, that is absorbed by a thick, dense, plasma is given by

$$f_a=\frac{4\pi}{cE^2}\int_0^\infty dx(j_{pe}\cdot E)_{\text{real}}$$
$$=\int_0^\infty dx\,\frac{4\pi e^2 n_{pe}\nu_{pe}}{m(\omega^2+\nu_{pe}^2)}\,e^{-2k_i x}\approx\frac{\nu_{pe}}{2k_i c}\,\frac{\omega_{pe}^2}{\omega^2}. \tag{A6}$$

The absorption reaches a maximum, $f_a=(\nu_{pe}/2\omega_{pe})^{1/2}$, at $\omega=\omega_{pe}$. But for $\omega<0.9\omega_{pe}$,

$$f_a<\frac{3}{2}\,\frac{\nu_{pe}}{\omega_{pe}}. \tag{A7}$$

In general, the plasma density does not fall abruptly to zero at the edges, and thus additional absorption occurs in the underdense plasma wings, where $\omega_{pe}(x)\leqslant\omega$. In that region power absorbed is absorbed at rate $2k_i(x)$, where $k_i\approx\nu_{pe}\omega_{pe}^2/2\omega c(\omega^2-\omega_{pe}^2)^{1/2}$ according to Eq. (A5). If n_{pe} falls linearly to zero over a distance Δx, the fraction of power absorbed in a round trip through the underdense wings is given at normal incidence by

$$\Delta f_a\approx 4\int_0^{\Delta x} dx k_i(x)$$
$$\approx\frac{2\nu_{pe}}{c}\int_0^{\Delta x} dx\,\frac{x}{\Delta x}\left(1-\frac{x}{\Delta x}\right)^{-1/2}=\frac{8}{3}\,\frac{\nu_{pe}\Delta x}{c}. \tag{A8}$$

The sum, $f_a+\Delta f_a$, gives the total fraction of power absorbed by the plasma.

Downloaded 18 Dec 2006 to 132.250.158.37. Redistribution subject to AIP license or copyright, see http://pop.aip.org/pop/copyright.jsp

APPENDIX B: TWO-STREAM INSTABILITY

As shown by Briggs,[11] the dispersion relationship for longitudinal waves $e^{i(kz-\omega t)}$ produced by a uniform, collisionless electron beam propagating through a uniform plasma is given by

$$1-\frac{\omega_{pe}^2}{\omega(\omega+i\nu_{pe})}-\frac{\omega_b^2}{k^2}\int_0^\infty dv_b\,\frac{1}{v_b-\omega/k}\,\frac{df}{dv_b}=0, \quad \text{(B1)}$$

where ν_{pe} is the plasma collision frequency, ω_{pe} and ω_b are the plasma frequencies for the plasma and beam, respectively, and $f(v_b)$ is the axial beam velocity distribution. To analyze Eq. (B1), choose k real and set $\omega=\omega_r+i\omega_i$. Since instability occurs only if $\omega_i>0$ over some range of k, consider a range where $\omega_{pe}\gg\omega_i,\ \nu_{pe},\ \omega_b$. The imaginary part of Eq. (B1) then reduces to

$$(2\omega_i+\nu_{pe})\,\frac{\omega_{pe}^2}{\omega_r^3}\approx\pi\,\frac{\omega_b^2}{k^2}\,\frac{df}{dv_b}, \quad \text{(B2)}$$

while the real part yields the resonance condition

$$\omega_r\approx\omega_{pe}\approx kv_b. \quad \text{(B3)}$$

Imposing condition (B3) on Eq. (B2) shows that ω_i is never positive if

$$\nu_{pe}\geqslant\frac{\pi\omega_b^2}{\omega_{pe}}\,v_b^2\,\frac{df}{dv_b}. \quad \text{(B4)}$$

If this condition is met for all v_b, the beam is stable. For example, if $f(v_b)$ decreases linearly to zero over a range Δv from a maximum velocity v_0,

$$\frac{df}{dv_b}=\frac{2}{\Delta v^2}, \quad \text{(B5)}$$

between $(v_0-\Delta v)$ and v_0. Stability criterion (B4) then reduces to

$$\nu_{pe}\geqslant\frac{2\pi\omega_b^2}{\omega_{pe}}\left(\frac{v_0}{\Delta v}\right)^2=2\pi\left(\frac{n_b}{n_{pe}}\right)\left(\frac{v_0}{\Delta v}\right)^2\omega_{pe}. \quad \text{(B6)}$$

[1] W. M. Manheimer, IEEE Trans. Plasma Sci. **PS-19**, 1228 (1991).
[2] M. A. Lieberman and A. J. Lichtenberg, *Principles of Plasma Discharges and Materials Processing* (Wiley, New York, 1994), Chap. 1, 3, and 14.
[3] J. Mathew, R. F. Fernsler, R. A. Meger, J. A. Gregor, D. P. Murphy, R. E. Pechacek, and W. M. Manheimer, Phys. Rev. Lett. **77**, 1982 (1996).
[4] R. A. Meger, J. Mathew, J. A. Gregor, R. E. Pechacek, R. F. Fernsler, W. M. Manheimer, and A. E. Robson, Phys. Plasmas **2**, 2532 (1995).
[5] A. E. Robson, R. L. Morgan, and R. A. Meger, IEEE Trans. Plasma Sci. **PS-20**, 1036 (1992).
[6] R. D. Hake, Jr. and A. V. Phelps, Phys. Rev. **158**, 70 (1967).
[7] J. D. Jackson, *Classical Electrodynamics* (Wiley, New York, 1967), Chap. 13.
[8] T. Majeed and D. J. Strickland, J. Phys. Chem. Ref. Data **26**, 335 (1997).
[9] L. G. Christophorou, *Atomic and Molecular Radiation Physics* (Wiley, New York, 1971), pp. 34–40.
[10] F. J. Mehr and M. A. Biondi, Phys. Rev. **181**, 264 (1969).
[11] R. J. Briggs, Adv. Plasma Phys. **4**, 43 (1971).
[12] Published values of γ vary widely. The values listed in Ref. 2 show $\gamma<0.3$, but values above unity have been reported; see, for example, B. Szapier, J. J. Rocca, and T. Prabhuram, Appl. Phys. Lett. **53**, 358 (1988).
[13] P. F. Little and A. von Engel, Proc. R. Soc. London, Ser. A **224**, 209 (1954).
[14] M. M. Kuraica and N. Konjevic, Appl. Phys. Lett. **70**, 1521 (1997).
[15] I. Langmuir, Phys. Rev. **33**, 954 (1929).
[16] D. Rapp and P. Englander-Golden, J. Chem. Phys. **43**, 1464 (1965).
[17] V. I. Kolobov and L. D. Tsendin, Plasma Sources Sci. Technol. **4**, 551 (1995).
[18] R. F. Fernsler, Bull. Am. Phys. Soc. **42**, 2006 (1997).
[19] R. F. Fernsler, R. F. Hubbard, and M. Lampe, J. Appl. Phys. **75**, 3278 (1994).
[20] M. Hayashi, ''Recommended values of transport cross sections for elastic collisions and total cross section for electrons in atomic and molecular gases,'' NTIS Document No. N82-18007, 1981. Copies may be ordered through the National Technical Information Service, Springfield, VA 22161.

Downloaded 18 Dec 2006 to 132.250.158.37. Redistribution subject to AIP license or copyright, see http://pop.aip.org/pop/copyright.jsp